

The rho meson in a scenario of pure chiral restoration

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Abstract

Based on QCD sum rules we explore the consequences of a pure chiral restoration scenario for the ρ meson, where all chiral symmetry breaking condensates are dropped whereas the chirally symmetric condensates remain at their vacuum values. This pure chiral restoration scenario causes the drop of the ρ spectral moment by about 120 MeV. The complementarity of mass shift and broadening is discussed. A simple parametrization of the ρ spectral function leads to a width of about 600 MeV if no shift of the peak position is assumed.

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1. Introduction: The impact of chiral symmetry restoration on the properties of hadrons is a much debated issue. In particular light vector mesons have been studied extensively both on the theoretical and the experimental side; for recent reviews see e.g. [1–5]. In fact, in-medium modifications of hadrons made out of light quarks and especially their possible ”mass drops” are taken often synonymously for chiral restoration. The Brown-Rho scaling conjecture [6] and Ioffe’s formula for the nucleon [7] suggest such a tight connection. However, experimentally the main observation of in-medium changes of light vector mesons via dilepton spectra is a significant broadening of the spectral shape [8, 9]. Such a broadening can be obtained in hadronic many-body approaches, e.g., [10–14], which at first sight are not related directly to chiral restoration in the above spirit. Pion dynamics and resonance formation, both fixed to vacuum data, provide the important input for such many-body calculations. Clearly the pion dynamics is closely linked to the vacuum phenomenon of chiral symmetry *breaking*, but the connection to chiral *restoration* is not so clear. For the physics of resonances the connection is even more loose. There are recent approaches which explain some hadronic resonances as dynamically generated from chiral dynamics [15–21], but again this primarily points towards an intimate connection between hadron physics and chiral symmetry breaking and not so much chiral restoration. As suggested, e.g., in [2, 22] the link to chiral restoration might be indirect: The in-medium broadening could be understood as a step towards deconfinement. In the deconfined quark-gluon plasma also chiral symmetry is presumed to be restored. All these considerations suggest that the link between chiral restoration and in-medium changes of hadrons is not as clear as one might have hoped.

Additional input could come from approaches which are closer to QCD than standard hadronic models. One such approach is the QCD sum rule method [23–27]. A somewhat superficial view on QCD sum rules for vector mesons seems to support the original picture of an intimate connection between chiral restoration and in-medium changes. Here the previously popular chain of arguments goes as follows: (1) Four-quark condensates play an important role for the vacuum mass of the light vector mesons [23, 24] (see, however, [28] for a different view). (2) The four-quark condensates factorize into squares of the two-quark condensate [29]. (3) The two-quark condensate drops in the medium due to chiral restoration [30, 31]. (4) Thus the four-quark condensates drop in the medium accordingly. (5) Therefore the masses of light vector mesons change (drop) in the medium due to chiral restoration.

In this line of reasoning only the points 1 and 3 are undoubted. Even if one follows the

arguments of points 1 to 4 it has been shown that besides a dropping mass also a broadened hadronic spectral distribution is compatible with the QCD sum rules [32, 33]. One still seems to have at least a connection between chiral restoration — drop of two- and four-quark condensates — and in-medium changes, no matter whether it is a mass shift or a broadening or a more complicated in-medium modification [13, 34]. However, also point 2 and as its consequence point 4 are questionable: Whether the four-quark condensates factorize at least in vacuum is discussed since the invention of QCD sum rules, see, e.g., [23, 24, 35–41] and for in-medium situations also [27, 42–46]. Raising doubts on point 2 immediately questions point 4 and in that way the seemingly clear connection between chiral restoration and in-medium changes gets lost.

Indeed, a closer look on the sum rules for light vector mesons reveals that most of the condensates, whose in-medium change is translated into an in-medium modification for the respective hadron, are actually chirally symmetric (see below). Physically, it is of course possible and by far not unreasonable that the same microscopic mechanism which causes the restoration of chiral symmetry is also responsible for changes of chirally symmetric condensates. For example, in the scenario [47] about half of the (chirally symmetric) gluon condensate vanishes together with the two-quark condensate. On the other hand, these considerations show that the connection between the mass of a light vector meson and chiral symmetry breaking is not as direct as one would naively expect.

We take these considerations as a motivation to study in the present work a “pure chiral restoration” scenario, i.e. to ask the question: How large would the mass or the width of the ρ meson be in a world where only chiral symmetry breaking objects/condensates are dropped. We stress that such a scenario may not reflect all the physics which is contained in QCD. There might be intricate interrelations between chirally symmetric and symmetry breaking objects. In that sense the pure chiral restoration scenario shows the minimal impact that the restoration of chiral symmetry has on the properties of the ρ meson.

2. Chiral transformations and QCD condensates: For vanishing quark masses, QCD with N_f flavors is invariant with respect to the global chiral $SU_R(N_f) \times SU_L(N_f)$ transformations. Focusing for the time being on the $N_f = 2$ light (massless) quark sector, the corresponding left-handed transformations read for the left-handed quark field $\psi_L =$

$\frac{1}{2}(1 - \gamma_5)\psi$ and the right-handed quark field $\psi_R = \frac{1}{2}(1 + \gamma_5)\psi$

$$\psi_L \rightarrow e^{i\vec{\theta}_L \cdot \vec{\tau}} \psi_L, \quad \psi_R \rightarrow \psi_R, \quad (1)$$

while the right-handed transformations are

$$\psi_R \rightarrow e^{i\vec{\theta}_R \cdot \vec{\tau}} \psi_R, \quad \psi_L \rightarrow \psi_L, \quad (2)$$

where $\vec{\tau}$ are the iso-spin Pauli matrices and $\psi = \begin{pmatrix} u \\ d \end{pmatrix}$ denotes the quark iso-doublet. Equations (1, 2) represent isospin transformations acting separately on the right-handed and left-handed parts of the quark field operator $\psi = \psi_L + \psi_R$, i.e. the three-component vectors $\vec{\theta}_{R,L}$ contain arbitrary real numbers. Gluons and heavier quarks remain unchanged with respect to the transformations (1, 2).

A quark current which has the quantum numbers of the ρ meson is given by the vector-iso-vector current

$$\vec{j}^\mu = \frac{1}{2} \bar{\psi} \gamma^\mu \vec{\tau} \psi. \quad (3)$$

If a chiral transformation according to (1, 2) is applied to \vec{j}^μ , it becomes mixed with the axial-vector-iso-vector current

$$\vec{j}_5^\mu = \frac{1}{2} \bar{\psi} \gamma^\mu \gamma_5 \vec{\tau} \psi \quad (4)$$

which carries the quantum numbers of the a_1 meson. Indeed, experiments show that the vector current (3) couples strongly to the ρ meson, while the axial-vector current (4) couples to the a_1 meson [48]. Therefore, ρ and a_1 are called chiral partners.

The central object of QCD sum rules [23, 24, 49] is the retarded current-current correlator which reads for the ρ^0 meson

$$\Pi^{\mu\nu}(q) = i \int d^4x e^{iqx} \Theta(x_0) \langle [j_3^\mu(x), j_3^\nu(0)] \rangle, \quad (5)$$

where for vacuum ($\langle \dots \rangle$ means accordingly the vacuum expectation value) the retarded and time-ordered propagator coincide for positive energies, whereas for in-medium situations (e.g. nuclear matter, $\langle \dots \rangle$ refers then to the Gibbs average), the retarded correlator has to be taken (cf. [27]). The imaginary part of the current-current correlator contains the spectral distribution, i.e. the information which hadronic one-body and many-body states couple to the considered current. For large space-like momenta, $Q^2 \equiv -q^2 \gg \Lambda_{QCD}^2$, the correlator can be reliably calculated from the elementary QCD quark and gluon degrees of freedom due to

asymptotic freedom. Results from QCD perturbation theory can be systematically improved by the introduction of quark and gluon condensates using the operator-product expansion (OPE) [50]. The QCD sum rule method connects the mentioned two representations of the correlator by a dispersion relation which reads after a Borel transformation

$$\frac{1}{\pi} \int_0^\infty ds s^{-1} \text{Im} \Pi(s) e^{-s/M^2} = \tilde{\Pi}(M^2), \quad (6)$$

where the Borel mass M has emerged from the OPE momentum scale Q (for further details we refer the interested reader to [33]). We consider a ρ meson at rest, therefore, the tensor structure of (5) reduces to a scalar $\Pi = \frac{1}{3}\Pi_\mu^\mu$. The Borel-transformed OPE reads

$$\tilde{\Pi}(M^2) = c_0 M^2 + \sum_{i=1}^{\infty} \frac{c_i}{(i-1)! M^{2(i-1)}} \quad (7)$$

with coefficients up to mass dimension 6

$$c_0 = \frac{1}{8\pi^2} \left(1 + \frac{\alpha_s}{\pi}\right), \quad (8)$$

$$c_1 = -\frac{3}{8\pi^2} (m_u^2 + m_d^2), \quad (9)$$

$$c_2 = \frac{1}{2} \left(1 + \frac{\alpha_s}{4\pi} C_F\right) (m_u \langle \bar{u}u \rangle + m_d \langle \bar{d}d \rangle) + \frac{1}{24} \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle + N_2, \quad (10)$$

$$c_3 = -\frac{112}{81} \pi \alpha_s \langle \mathcal{O}_4^V \rangle - 4N_4 \quad (11)$$

with $C_F = (n_c^2 - 1)/(2n_c) = 4/3$ for $n_c = 3$ colors. A mass dimension 2 condensate seems to be excluded in vacuum [51]. In (8 - 11) we have introduced the strong coupling α_s , the light-quark masses $m_{u,d}$, the gluon condensate $\langle \frac{\alpha_s}{\pi} G^2 \rangle$, and the combination of four-quark condensates in compact notation

$$\langle \mathcal{O}_4^V \rangle = \frac{81}{224} \langle (\bar{\psi} \gamma_\mu \gamma_5 \lambda^a \tau_3 \psi)^2 \rangle + \frac{9}{112} \langle \bar{\psi} \gamma_\mu \lambda^a \psi \sum_{f=u,d,s} \bar{f} \gamma^\mu \lambda^a f \rangle \quad (12)$$

with color matrices λ^a . (For a classification of four-quark condensates cf. [52]. Here, we have extended the notation to the SU(3) flavor sector.) It is also useful to introduce the averaged two-quark condensate $m_q \langle \bar{q}q \rangle = \frac{1}{2} \langle m_u \bar{u}u + m_d \bar{d}d \rangle$ and $m_q = (m_u + m_d)/2$. These terms constitute the contributions which already exist in vacuum (and might change in a medium) up to higher-order condensates (including, for instance, the poorly known term c_4) which are suppressed by higher powers in the expansion parameter M^{-2} . Additional non-scalar condensates come into play, in particular for in-medium situations. In (10, 11), only the twist-two non-scalar condensates [27] are displayed, $N_i = -\frac{2}{3} i \langle \mathcal{ST} \bar{\psi} \gamma_{\mu_1} D_{\mu_2} \dots D_{\mu_i} \psi \rangle g^{\mu_1 0} \dots g^{\mu_i 0}$,

where the operation \mathcal{ST} is introduced to make the operators symmetric and traceless with respect to its Lorentz indices. Twist-four non-scalar condensates have been found to be numerically less important [27, 53, 54].

Using (1, 2) one can show that the only objects in the OPE (7) with coefficients (8 - 11) which are *not* chirally invariant¹ are (i) the (numerically small) two-quark condensate and (ii) a part of the (numerically important) four-quark condensate $\langle \mathcal{O}_4^V \rangle$ specified in (12). One can split the four-quark condensates (12) into a chirally symmetric part,

$$\langle \mathcal{O}_4^{\text{sym}} \rangle = \frac{81}{448} \langle (\bar{\psi} \gamma_\mu \gamma_5 \lambda^a \tau_3 \psi)^2 + (\bar{\psi} \gamma_\mu \lambda^a \tau_3 \psi)^2 \rangle + \frac{9}{112} \langle \bar{\psi} \gamma_\mu \lambda^a \psi \sum_{f=u,d,s} \bar{f} \gamma^\mu \lambda^a f \rangle, \quad (13)$$

and a part which can be transformed into its negative by a proper chiral transformation (dubbed “chirally odd” object),

$$\langle \mathcal{O}_4^{\text{br}} \rangle = -\frac{81}{112} \langle (\bar{\psi}_R \gamma_\mu \lambda^a \tau_3 \psi_R) (\bar{\psi}_L \gamma^\mu \lambda^a \tau_3 \psi_L) \rangle \quad (14)$$

with $\langle \mathcal{O}_4^V \rangle = \langle \mathcal{O}_4^{\text{sym}} \rangle + \langle \mathcal{O}_4^{\text{br}} \rangle$. The last term in (13) is an iso-singlet. In an isospin invariant system, the other terms in (13) may be written as $\langle (\bar{\psi} \gamma_\mu \gamma_5 \vec{\tau} \lambda^a \psi)^2 + (\bar{\psi} \gamma_\mu \vec{\tau} \lambda^a \psi)^2 \rangle = 2 \langle (\bar{\psi}_R \gamma_\mu \vec{\tau} \lambda^a \psi_R)^2 + (\bar{\psi}_L \gamma_\mu \vec{\tau} \lambda^a \psi_L)^2 \rangle$. The latter two terms are separately invariant with respect to left-handed and right-handed isospin transformations, i.e. they are chirally invariant.

3. Pure chiral restoration scenario for the rho meson: It appears to be very natural that a four-quark condensate which breaks chiral symmetry can be related to the square of the two-quark condensate which also breaks chiral symmetry. Indeed, it has been shown in [41] that the factorization of the four-quark condensate $\langle \mathcal{O}_4^{\text{br}} \rangle$ given in (14) is completely compatible with the ALEPH data on the vector and axial-vector spectral distributions [48]. On the other hand, it is not so obvious that a chirally symmetric four-quark condensate like $\langle \mathcal{O}_4^{\text{sym}} \rangle$ in (13) is related directly to the two-quark condensate.

Along this line of arguments we are going to answer the following question. What happens to the ρ meson if one keeps all expectation values of chirally invariant operators at their vacuum values and puts all chiral symmetry breaking objects to zero? This question defines what is meant by the “pure chiral restoration” scenario.

Let us first describe briefly how we fix the numerical values for the QCD condensates, collected in Tab. I. We note that the non-scalar condensates, which appear in (8 - 11),

¹ Note that c_1 breaks the chiral symmetry explicitly. Its contribution is numerically completely negligible.

QCD condensate	transformation	vacuum value
$\langle \bar{q}q \rangle$	chirally odd	$-(240 \text{ MeV})^3$
$\langle \frac{\alpha_s}{\pi} G^2 \rangle$	invariant	$(330 \text{ MeV})^4$
$\langle \mathcal{O}_4^{\text{sym}} \rangle$	invariant	$(267 \text{ MeV})^6$
$\langle \mathcal{O}_4^{\text{br}} \rangle$	chirally odd	$\frac{9}{7} \langle \bar{q}q \rangle^2$

TABLE I: Employed QCD condensates, their behavior with respect to chiral transformations and their respective size.

are chirally symmetric and vanish in the vacuum, i.e. we can disregard them also for the scenario of pure chiral restoration. Next we turn to the vacuum condensates. The gluon condensate is determined from the QCD sum rules for the charmonium [23, 24]. The running coupling has to be evaluated at the scale M . Following [55] we use $\alpha_s = 0.38$. The two-quark condensate is fixed by the Gell-Mann–Oakes–Renner relation [56] $m_q \langle \bar{q}q \rangle = -\frac{1}{2} F_\pi^2 M_\pi^2$ with the pion-decay constant $F_\pi \approx 92 \text{ MeV}$, $m_q = (m_u + m_d)/2$ and the pion mass $M_\pi \approx 140 \text{ MeV}$ [57]. Using in addition $m_q = 6 \text{ MeV}$ [58] one gets $\langle \bar{q}q \rangle = -(240 \text{ MeV})^3$. For vacuum, the condensate $\langle \mathcal{O}_4^{\text{br}} \rangle$ has been extracted from the experimental difference between vector and axial-vector spectral information [48]. We use the result of [41]

$$\langle \mathcal{O}_4^{\text{br}} \rangle_{\text{vac}} \approx \frac{9}{7} \langle \bar{q}q \rangle_{\text{vac}}^2 \quad (15)$$

together with the vacuum ρ -meson properties [57] to fix the vacuum value for $\langle \mathcal{O}_4^{\text{sym}} \rangle$ defined in (13). For the pure chiral restoration scenario one drops then the chiral symmetry breaking terms, in particular (14). Thus also here one only needs the *vacuum* values for the chirally invariant terms, in particular for (13). We recall that (15) indicates that the chiral symmetry breaking four-quark condensate factorizes [41].

To describe the properties of the ρ meson we rearrange (6) by splitting the integral $\int_0^\infty = \int_0^{s_+} + \int_{s_+}^\infty$ and putting the so-called continuum part to the OPE terms thus isolating the interesting hadronic resonance part below the continuum threshold s_+ . This allows to define the normalized moment [59] of the hadronic spectral function (16)

$$\tilde{m}^2(M, s_+) \equiv \frac{\int_0^{s_+} ds \text{Im}\Pi(s) e^{-s/M^2}}{\int_0^{s_+} ds \text{Im}\Pi(s) s^{-1} e^{-s/M^2}} \quad (16)$$

$$= \frac{c_0 M^2 [1 - (1 + \frac{s_+}{M^2}) e^{-s_+/M^2}] - \frac{c_2}{M^2} - \frac{c_3}{M^4} - \frac{c_4}{2M^6}}{c_0 [1 - e^{-s_+/M^2}] + \frac{c_1}{M^2} + \frac{c_2}{M^4} + \frac{c_3}{2M^6} + \frac{c_4}{6M^8}}, \quad (17)$$

where the semi-local duality hypothesis $\frac{1}{\pi} \int_{s_+}^{\infty} ds s^{-1} \text{Im}\Pi(s) e^{-s/M^2} = \int_{s_+}^{\infty} ds c_0 e^{-s/M^2}$ is exploited in (17). The second line emerges essentially from the OPE and "measures" how \tilde{m} is determined by condensates.

The meaning of the spectral moment (16) becomes obvious for the pole ansatz [24] of the hadronic spectral function below s_+ , $\text{Im}\Pi(s) = F_0 \delta(s - m_0^2)$, where $\tilde{m} = m_0$ follows and F_0 is determined by inserting m_0 in (6). For the sake of clarity let us consider first the vacuum case where we identify the average of \tilde{m} with the vacuum mass. The averaged mass parameter is determined by $\overline{m}(s_+) = (M_{\max} - M_{\min})^{-1} \int_{M_{\min}}^{M_{\max}} \tilde{m}(M, s_+) dM$ within the Borel window. According to [60] the Borel minimum is determined by the requirement that the mass dimension 6 contribution to the OPE is smaller than 10%. For the Borel maximum we demand that the continuum contribution to the spectral integral is smaller than 50%. s_+ follows from the requirement of maximum flatness of $\tilde{m}(M, s_+)$ as a function of M within the Borel window. Employing this system of equations, $\langle \mathcal{O}_4^{\text{sym}} \rangle = (267 \text{ MeV})^6$ with $\langle \mathcal{O}_4^{\text{br}} \rangle$ from (15) (for the other condensates see Tab. 1) is required to get $\overline{m} = 775.5 \text{ MeV}$ for $s_+ = 1.37 \text{ GeV}^2$, see upper curves in Fig. 1.

Let us consider now the pure chiral restoration scenario. $\langle \mathcal{O}_4^{\text{br}} \rangle \rightarrow 0$ and $\langle \bar{q}q \rangle \rightarrow 0$ but keeping the other condensate values causes a drop of \overline{m} to 659.8 MeV and the continuum threshold becomes $s_+ = 1.03 \text{ GeV}^2$, see lower curves in Fig. 1. We emphasize the large impact of dropping $\langle \mathcal{O}_4^{\text{br}} \rangle$ on the averaged spectral moment \overline{m} .

To get an estimate of the possible importance of the poorly known term c_4 we use as "natural scale" $\langle \frac{\alpha_s}{\pi} G^2 \rangle^2$. The Borel curves for $|c_4| = \langle \frac{\alpha_s}{\pi} G^2 \rangle^2$ border the bands in Fig. 1. This estimate is quite rough as c_4 may contain also chirally odd condensates, whose drop is not accounted for in the pure chiral restoration scenario.

Summarizing the outcome of this numerical study, the pure chiral restoration scenario is characterized by a drop of the model-independent spectral moment \overline{m} by about 120 MeV.

4. Mass shift vs. broadening: While for a narrow resonance in vacuum the often employed pole + continuum ansatz is reasonable, the spectral distribution may get a more complex structure in a medium [10–14, 28, 32, 34, 59, 61]. In particular, one cannot decide, within the employed framework of QCD sum rules, whether a drop of \tilde{m} , and consequently \overline{m} , means a mass shift or a broadening or both. To make it explicit, we use a Breit-Wigner

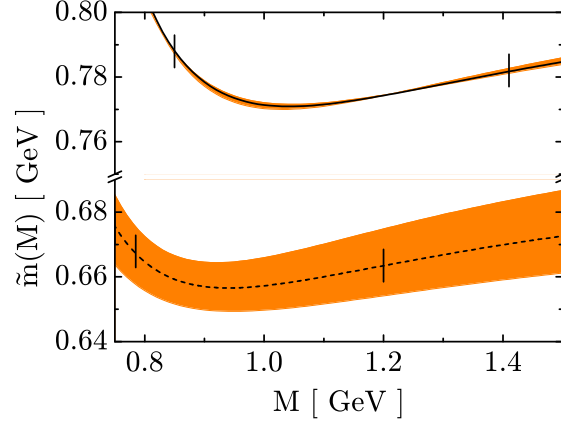


FIG. 1: The mass parameter $\tilde{m}(M, s_+)$ for the optimized continuum threshold s_+ as a function of the Borel mass M for vacuum values of condensates from Tab. I (upper curves) and the pure chiral restoration scenario with $\langle \mathcal{O}_4^{\text{br}} \rangle \rightarrow 0$ and $\langle \bar{q}q \rangle \rightarrow 0$ (lower curves). The Borel windows are marked by vertical bars. The curves are for $c_4 = 0$, while the bands cover the range $c_4 = \pm \langle \frac{\alpha_s}{\pi} G^2 \rangle^2$.

ansatz for the spectral function

$$\text{Im}\Pi(s \leq s_+) = \frac{F_0}{\pi} \frac{\sqrt{s}\Gamma(s)}{(s - m_0^2)^2 + s\Gamma^2(s)}, \quad (18)$$

where the vacuum parametrization of the width is given by [33]

$$\Gamma(s) = \Theta(s - 4m_\pi^2)\Gamma_0 \left(1 - \frac{4m_\pi^2}{s}\right)^{\frac{3}{2}} \left(1 - \frac{4m_\pi^2}{m_0^2}\right)^{-\frac{3}{2}}, \quad (19)$$

with $m_\pi = 135$ MeV being the pion mass. As a consequence of the two-parameter ansatz for the spectral function, the moment $\tilde{m}^2(M)$ determines a relation $m_0(M) = m_0(M, \Gamma_0)$; $F_0(M)$ is again determined by (6), and hence $\overline{m_0} = \overline{m_0}(\Gamma_0)$.

Adjusting $\langle \mathcal{O}_4^{\text{sym}} \rangle$ to reproduce the experimental mass $\overline{m_0} = 775.5$ MeV and width $\Gamma_0 = 149.4$ MeV we obtain now $\langle \mathcal{O}_4^{\text{sym}} \rangle = (242 \text{ MeV})^6$.

In general, the peak position m_{peak} and the full width at half maximum Γ_{FWHM} of the spectral function do not coincide with the corresponding parameters m_0 (or $\overline{m_0}$) and Γ_0 of the ansatz (18). For $\Gamma(s) = \Gamma_0 = \text{const.}$, e.g., m_0 is determined by $m_0^2 = \sqrt{4m_{\text{peak}}^4 + \Gamma_0^2 m_{\text{peak}}^2} - m_{\text{peak}}^2$. While for small Γ_0 the peak position m_{peak} and m_0 differ only by a few MeV, they differ significantly for larger values of Γ_0 . Especially, keeping the parameter $\overline{m_0}$ constant in the chiral symmetry restoration scenario causes an implicit shift of the peak position m_{peak} due to the broadening caused by the symmetry restoration.

Therefore, instead of requiring one of the possible options $\overline{m_0}$ or Γ_0 to be constant, we now demand that the associated shape characteristics of the spectral function, i.e. m_{peak}

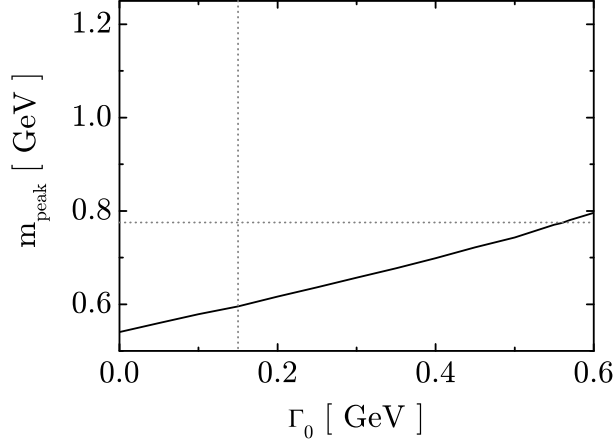


FIG. 2: Peak position m_{peak} as a function of the width parameter Γ_0 for $\langle \mathcal{O}_4^{\text{sym}} \rangle = (242 \text{ MeV})^6$. Dotted lines mark the experimental values $m_{\text{peak}} = 775.5 \text{ MeV}$ and $\Gamma_{\text{FWHM}} = 149.4 \text{ MeV}$.

or Γ_{FWHM} , are fixed within the restoration procedure. In doing so, it is convenient to set $\Gamma(s) = \Gamma_0 = \text{const}$ for the chiral symmetry restoration scenario. The result is depicted in Fig. 2. For the pure chiral restoration scenario, the curve $m_{\text{peak}}(\Gamma_0)$ is significantly shifted away from the vacuum physical point $(m_0, \Gamma_0) = (775.5 \text{ MeV}, 149.4 \text{ MeV})$. If one assumed that chiral restoration in the present spirit does not cause an additional broadening, one would recover the previously often anticipated "mass drop".

Fig. 2 evidences, however, that an opposite interpretation is conceivable as well, namely pure broadening with keeping the vacuum value of m_{peak} . The NA60 [8, 9] and CLAS [62] data seem indeed to favor such a broadening effect. In fact, assuming that m_{peak} does not change by chiral restoration, the width is increased to 600 MeV. In this respect, broadening of a spectral function signals equally well chiral restoration as dropping mass would do.

Fig. 3 exhibits the spectral function $\text{Im}\Pi(s)$ as a function of s for the two extreme options above. The solid curve depicts the enormous broadening when keeping the peak at the vacuum position. The dashed curve is for the dropping mass option for keeping the full width at half maximum at its vacuum value. From the perspective of the employed QCD sum rules, both options are equivalent, as any other point on the curve $m_{\text{peak}}(\Gamma_0)$ in Fig. 2.

The overall outcome seems to be that the Borel transformed QCD sum rule requires more strength of the spectral function at lower energies. This may be realized by a shift or a broadening or both.

It should be emphasized, in this context, that (18) is a pure ad hoc ansatz. For in-

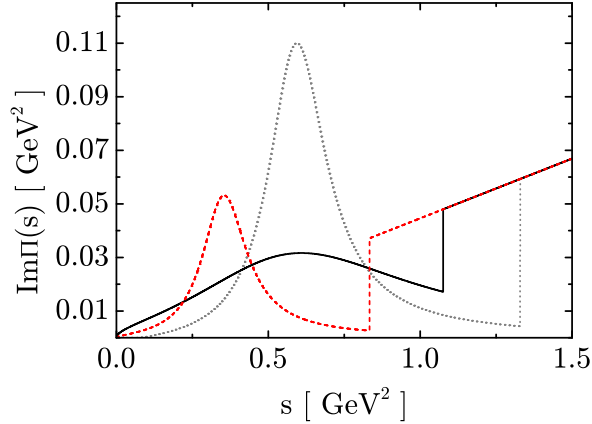


FIG. 3: The spectral density $\text{Im}\Pi(s)$ in the vacuum case (dotted curve) and in the pure chiral restoration scenario for $\Gamma_{\text{FWHM}} = \text{const}$ (dashed curve) or $m_{\text{peak}} = \text{const}$ (solid curve).

stance, multi-peak structures may emerge from particle-hole and resonance-hole excitations in the nuclear medium [11–13, 61, 63]. It is desirable, therefore, to determine the spectral distribution by experimental data of ρ decay into dileptons in a nuclear medium to decide whether it deviates from the vacuum and try to relate it to the change of condensates and Landau term. Such strategy for a particular model instead of data is also envisaged in [28] in a truncated hierarchy of spectral moments. Clearly a better microscopic understanding is mandatory for an appropriate modelling of the spectral shape. This applies also for the continuum threshold region.

5. Notes on omega and axial-vector mesons: The current $\bar{u}\gamma^\mu u + \bar{d}\gamma^\mu d$ has the quantum numbers of the ω meson and is a chiral singlet with respect to $SU_R(2) \times SU_L(2)$ chiral transformations. Consequently, the ω meson does not have a chiral partner *in a world with two light flavors*. Concerning *three* light flavors the current given above is a superposition of a member of the flavor octet and the singlet. The octet members do have chiral partners and one may assign a proper linear combination of the two f_1 [57] mesons as the chiral partner of the respective linear combination of ω and ϕ .

All considerations made in the following concern *two* light flavors. The OPE side of the ω meson including terms up to dimension 6 contains only chirally symmetric terms [23, 24, 27] – except for the two-quark condensate term $\propto m_q \langle \bar{q}q \rangle$. This term, however, breaks chiral symmetry explicitly by the quark mass and dynamically by the quark condensate. Considerations about symmetry transformations, on the other hand, concern the case where chiral symmetry is exact, i.e. without explicit breaking. Therefore the appearance of the

term $\propto m_q \langle \bar{q}q \rangle$ is not in contradiction to the statement that the ω is a chiral singlet. Without explicit calculations it is clear that in the pure chiral restoration scenario the ω meson does not change much of its mass since only the numerically very small term $\propto m_q \langle \bar{q}q \rangle$ is dropped while the chirally symmetric four-quark condensates do not change.

The current (4) with the quantum numbers of the a_1 meson yields the same chirally symmetric OPE parts as the ρ meson. The chirally odd parts are of course different, they are the negative of the ones which appear in the OPE for the ρ meson [23, 24, 27]. In addition, the hadronic side of the sum rule contains not only the a_1 , but also the pion. The latter contribution is $\propto F_\pi^2$ where F_π denotes the pion-decay constant. Both effects, different chirally odd condensates and the appearance of the pion, lead to the fact that the sum rule method yields a mass for the a_1 which is significantly different from the ρ meson mass [23, 24] – as it should be. In the pure chiral restoration scenario the chirally odd condensates are put to zero. In addition, the pion-decay constant which is an order parameter of chiral symmetry breaking [64] also vanishes. Then the sum rules for ρ and a_1 are the same. As expected the chiral partners become degenerate in the pure chiral restoration scenario.

6. Summary: Two extreme and antagonistic statements concerning hadron masses and hadronic in-medium modifications could be raised: (a) Basically all hadron masses are caused by chiral symmetry breaking. Consequently in a dense and/or hot strongly interacting medium the masses of hadrons vanish at the point of chiral restoration – apart from some small remainder which is due to the explicit breaking of chiral symmetry by the finite quark masses. (b) The observed in-medium changes can be explained by standard hadronic many-body approaches and have no direct relation to chiral restoration. Our findings do not support either of these extreme statements. If one drops the chiral symmetry breaking condensates in the sum rule for the ρ meson one does see a significant change of the mass moment. This is neither 100 % (as statement (a) would suggest) nor 0 % (statement (b)). In the scenario of pure chiral restoration we have kept the chirally invariant condensates at their vacuum values. Though, an adequate restoration mechanism might also change these condensates, this allows a discussion of chiral-symmetry restoration which is not interfered by additional in-medium effects. This is clearly unrealistic for a true in-medium situation. In particular, the contributions coming from the non-scalar twist-two operators are found to be sizable, e.g., for cold nuclear matter [26]. Nonetheless, our scenario indicates that the connection between the vacuum masses and chiral symmetry breaking or between dropping

masses and chiral restoration is not direct. In principle, one could imagine conspiracies between chiral symmetry breaking and non-breaking condensates such that one of the extreme statements raised above becomes true. Such a conspiracy would be driven by the underlying microscopic mechanisms which cause spontaneous chiral symmetry breaking and/or its restoration. Clearly we need a deeper understanding of these microscopic mechanisms. Finally, we emphasize that, within the framework of QCD sum rules, a “dropping mass” and a broadening of the spectral function are both equally well conceivable in their relation to chiral restoration.

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